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NONLINEAR STABILIZATION OF THE CURRENT CONVECTIVE INSTABILITY I--ETC(U)
JAN 80 P K CHATURVEDI, S L OSSAKOW
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*Present address: Berkeley Research Associates, Arlington, Va. 22209.

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The current convective instability has recently been proposed as a possible mechanism for generating large scale irregularities in the auroral F region, driven by weak field-aligned currents associated with the diffuse aurora and a plasma density gradient [Ossakow and Chaturvedi, 1979]. In this report, we show that this instability can stabilize nonlinearly by generating linearly damped harmonics. An estimate of the saturated amplitudes show good agreement with the available experimental data. In addition, with a northward plasma density gradient, the dominant nonlinear harmonic is in the northward direction, in agreement with DNA Wideband satellite observations.

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NONLINEAR STABILIZATION OF THE CURRENT CONVECTIVE INSTABILITY IN THE DIFFUSE AURORA

The high latitude scintillation enhancements observed recently by the DNA Wideband satellite have been characterized by the association of their occurrence regions with horizontal plasma density gradients and with the nighttime diffuse auroral particle precipitation zone [*Fremouw et al.*, 1977; *Rino et al.*, 1978]. The scintillations are believed to be caused by field-aligned sheetlike F region ionospheric irregularities. Recently, we have argued that a current convective instability, driven by a weak magnetic field-aligned current and a plasma density gradient transverse to the magnetic field, could generate large scale-size irregularities in the medium; even when the directions of the ambient horizontal electric field and density gradient are such that the usual $\underline{E} \times \underline{B}$ gradient drift instability is inoperative [*Ossakow and Chaturvedi*, 1979]. It was shown there that the threshold conditions for the instability could readily be exceeded by the observed values of weak diffuse auroral currents and horizontal plasma density gradients at F region altitudes. In this report, we show that the current convective instability saturates nonlinearly by generating linearly damped harmonics through a dominant two-dimensional nonlinearity in the continuity equation. The estimated saturated amplitudes are then compared with the experimental data recently obtained through in-situ rocket flights (M. C. Kelley, private communication, 1979), and the agreement is found to be satisfactory. We note here that this nonlinear saturation mechanism has earlier been studied for the cases of some other gradient driven instabilities, namely, the $\underline{E} \times \underline{B}$ gradient-drift instability [*Rognlien and Weinstock*, 1974; *Chaturvedi and Ossakow*, 1979] and the collisional and collisionless gravitational (Rayleigh-Taylor) instability [*Chaturvedi and Ossakow*, 1977; *Hudson*, 1978].

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We consider a coordinate system appropriate for the auroral F region ionosphere in which the earth's magnetic field is aligned with the z-axis, the density gradient points northward along the y-axis, and the x-axis points westward. Note that an ambient electric field in the east-west direction and a horizontal gradient in the east-west direction have also been occasionally observed. But we shall ignore them for simplicity as we do the vertical altitude density gradient which is very weak. The basic equations, describing a nearly field-aligned current-convective mode in this situation, are

$$\underline{v}_e = \frac{cT_e}{eB_0}\frac{\nabla_\perp n \times \hat{z}}{n} + \frac{c\,\underline{E}_\perp \times \hat{z}}{B_0} - \frac{\nu_{ei}\,c_s^2}{\Omega_e\Omega_i}\frac{\nabla_\perp n}{n}$$
$$- \frac{e\underline{E}_z}{m\nu_{ei}} - \left[\frac{T_e}{m\nu_{ei}} + \frac{c_s^2}{\nu_{in}}\right]\frac{1}{n}\frac{\partial n}{\partial z}\,\hat{z} + v_0\,\hat{z} \tag{1}$$

$$\underline{v}_i = \frac{c\underline{E}_\perp \times \hat{z}}{B_0} + \frac{\nu_{in}}{\Omega_i}\frac{c\underline{E}_\perp}{B_0} - \frac{cT_i}{eB_0}\frac{\nabla_\perp n \times \hat{z}}{n} - \frac{\nu_{in}cT_i}{\Omega_i eB_0}\frac{\nabla_\perp n}{n}$$
$$- \frac{\nu_{ei}}{\Omega_e}\frac{c_s^2}{\Omega_i}\frac{\nabla_\perp n}{n} - \frac{c_s^2}{\nu_{in}}\frac{1}{n}\frac{\partial n}{\partial z}\,\hat{z} + V_0\,\hat{z} \tag{2}$$

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n_\alpha \underline{v}_\alpha) = 0 \tag{3}$$

and

$$\nabla \cdot \underline{J} = 0 \tag{4}$$

where we allow for perturbations having small variations along the magnetic field.

The first two equations describe the motion of electron and ion fluids, in which we have ignored the inertial effects for the two species on account of the low frequencies considered. Equation (3) is the continuity equation for species α, where subscript α stands for ions or electrons; and Eq. (4) results from quasi-neutrality, $n_e \approx n_i \approx n$. The notation used here is standard: c is the speed of light, n is the density, v the fluid velocity, T the temperature in energy units, $\perp$ denotes perpendicular to the ambient magnetic field $\underline{B}_0\hat{z}$, $\hat{z}$ is the unit vector along the z-axis, ν_{ei} is the Coulomb collision frequency, ν_{in} is the ion-neutral collision frequency, c_s^2 is the ion sound speed ($\equiv (T_e + T_i)/M$), M is the ion mass, m the electron mass, Ω is the cyclo-

tron frequency, and so on. We have neglected ν_{en}, the electron neutral collision frequency and have assumed $\nu_\alpha/\Omega_\alpha << 1$, appropriate for the F region.

In what follows we equate the weak diffuse auroral currents observed in the ionosphere to a relative drift between electrons and ions, v_d ($v_d = v_0 - V_0$), along the magnetic field. The ion and electron continuity equations may be written as, in the ion frame,

$$\begin{aligned}\frac{\partial n}{\partial t} + \frac{c\underline{E}_\perp \times \hat{z}}{B_0} \cdot \nabla_\perp n + \nabla_\perp \cdot \left[\frac{\nu_{in}}{\Omega_i}\, n\, \frac{c\underline{E}_\perp}{B_0}\right] \\ - \nabla_\perp \cdot \left[\frac{\nu_{in}}{\Omega_i}\, \frac{cT_i}{eB_0}\, \nabla_\perp n\right] \\ - \nabla_\perp \cdot \left[\frac{\nu_{ei}}{\Omega_e}\, \frac{c_s^2}{\Omega_i}\, \nabla_\perp n\right] - \frac{\partial}{\partial z}\left[\frac{c_s^2}{\nu_{in}}\, \frac{\partial n}{\partial z}\right] = 0\end{aligned} \tag{5}$$

and

$$\begin{aligned}\frac{\partial n}{\partial t} + v_d \frac{\partial n}{\partial z} + \frac{c\underline{E}_\perp \times \hat{z}}{B_0} \cdot \nabla_\perp n - \nabla_\perp \cdot \left[\frac{\nu_{ei}}{\Omega_e}\, \frac{c_s^2}{\Omega_i}\, \nabla_\perp n\right] \\ - \frac{\partial}{\partial z}\left[n\, \frac{eE_z}{m\nu_{ei}}\right] - \frac{\partial}{\partial z}\left\{\left[\frac{T_e}{m\nu_{ei}} + \frac{c_s^2}{\nu_{in}}\right]\frac{\partial n}{\partial z}\right\} = 0\end{aligned} \tag{6}$$

Equation (4) is nothing but the difference of Eqs. (5) and (6),

$$\begin{aligned}\nabla_\perp \cdot \left[n\, \frac{\nu_{in}}{\Omega_i}\, \frac{c\underline{E}_\perp}{B_0} - \frac{\nu_{in}}{\Omega_i}\, \frac{cT_i}{eB_0}\, \nabla_\perp n\right] \\ + \frac{\partial}{\partial z}\left[\frac{T_e}{m\nu_{ei}}\, \frac{\partial n}{\partial z} + n\frac{eE_z}{m\nu_{ei}}\right] = v_d\, \frac{\partial n}{\partial z}\end{aligned} \tag{7}$$

Any two of the Eqs. (5)-(7) provide the complete description of the current convective instability. In the linear approximation, a perturbation of the form $\propto \exp i(k_x x + k_y y + k_z z - \omega t)$, where $\omega \equiv \omega_r + i\gamma$, grows at a rate γ, given by,

$$\gamma = \frac{k_z v_d \dfrac{k_x}{k_\perp^2 L}\, \dfrac{\nu_{ei}}{\Omega_e}}{\left[\dfrac{k_z^2}{k_\perp^2} + \dfrac{\nu_{ei}\nu_{in}}{\Omega_e\Omega_i}\right]} - \frac{\nu_{ei}}{\Omega_e\Omega_i}\, c_s^2\, k_\perp^2 \tag{8}$$

$$-\frac{c_s^2 k_z^2}{\nu_{in}}\left[1+\frac{\dfrac{\nu_{in}^2}{\Omega_i^2}}{\dfrac{\nu_{ei}\nu_{in}}{\Omega_e\Omega_i}+\dfrac{k_z^2}{k_\perp^2}}\right]$$

where, $k_\perp^2 = k_x^2 + k_y^2$, $k_\perp^2 >> k_z^2$ and $L^{-1} = n_0^{-1}(\partial n_0/\partial y)$. This is the growth rate expression also derived in our earlier paper for the current convective instability [*Ossakow and Chaturvedi*, 1979] with three differences. In Eq. (8) we have ignored any equilibrium transverse electric fields, have included the diffusion damping and allowed for k_y. [Note that (8) is the same as (10) in *Ossakow and Chaturvedi* [1979] but differs from their (9). This is because in arriving at (1) and (2), in the present work, the relative motion in the electron-ion collision term in both the electron and ion momentum equations has been retained.]

The current convective instability is characteristic of collisional magnetoplasma configurations with a field-aligned current and a transverse density gradient, and was first discussed in connection with the positive column of laboratory gas discharges [*Kadomstev*, 1965]. A physical picture for this instability has been given in our earlier paper [*Ossakow and Chaturvedi*, 1979]. Briefly, the equilibrium electron drift along the magnetic field with respect to ions, v_d, causes polarization charges to develop across density fluctuations associated with the nearly perpendicular mode ($k_\perp >> k_z$) [see Eq. (7)]. The resulting $\underline{E}_\perp \times \underline{B}_0$ convection of the fluid along the equilibrium density gradient carries density depletions into increasing density regions and vice versa. As a consequence, the perturbation appears to grow with respect to the background [Eq. (5) or (6)]. The instability occurs when the growth term overcomes the damping due to ambipolar diffusion. This convective instability picture is all too familiar with the other gradient instabilities such as the gravitational and $\underline{E} \times \underline{B}$ instabilities [see, e.g., *Sudan et al.*, 1973].

In such cases, a dominant stabilization mechanism has been shown to be the two-dimensional mode coupling to linearly damped harmonics through $\underline{E}_\perp \times \underline{B}_0 \cdot \nabla_\perp n$ term in the

continuity equation [*Rognlien and Weinstock*, 1974; *Chaturvedi and Ossakow*, 1977; 1979]. This stabilization mechanism may be viewed as resulting from a nonlinear flux of particles or a quasi-linear modification of the equilibrium density gradient which acts to limit the growth of perturbations. It may be noted that the $\underline{E}_\perp \times \underline{B}_0 \cdot \nabla_\perp n$ term in the case of plane wave perturbations is vanishingly small, in which case some other nonlinearity may be the dominant one. In the present case, however, for highly field-aligned perturbations, such that $(\nu_{ei}/\Omega_e) >> (k_z/k_\perp)^2$ the two-dimensional nonlinearity in the continuity equation is the most important one. Thus in the ensuing analysis, we shall treat Eq. (7) linearly and retain only this nonlinearity in the continuity Eq. (6). We shall follow the approach outlined in *Rognlien and Weinstock* [1974] and *Chaturvedi and Ossakow* [1977;1979].

Consider a perturbation of the form

$$\left(\frac{\tilde{n}}{n_0}\right)_{1,1} = A_{1,1} \sin (k_x x + k_z z - \omega t) \cos k_y y \tag{9}$$

Then the linearlized Eq. (7) leads to

$$(\phi)_{1,1} = \frac{(k_z v_d)\dfrac{\Omega_i}{\nu_{in}}\dfrac{B_0}{c}}{\left(k_\perp^2 + \dfrac{\Omega_e \Omega_i}{\nu_{ei}\nu_{in}} k_z^2\right)} A_{1,1} \cdot \cos (k_x x + k_z z - \omega t) \cos k_y y \tag{10}$$

where we have ignored some small terms (i.e., the terms proportional to temperature; this is justified for the large parallel and perpendicular wavelengths we are considering). The $\tilde{\underline{E}}_\perp \times \underline{B}_0 \cdot \nabla_\perp \tilde{n}$ term in the continuity equation may be written as

$$\frac{\partial}{\partial y}(\phi)_{1,1} \frac{\partial}{\partial x}\left(\frac{\tilde{n}}{n_0}\right)_{1,1} - \frac{\partial (\phi)_{1,1}}{\partial x}\frac{\partial}{\partial y}\left(\frac{\tilde{n}}{n_0}\right)_{1,1} = -\frac{\dfrac{B_0}{c}\dfrac{\nu_{ei}}{\Omega_e}\dfrac{k_x k_y}{k_\perp^2} k_z v_d}{\left(\dfrac{k_z^2}{k_\perp^2} + \dfrac{\nu_{ei}\nu_{in}}{\Omega_e \Omega_i}\right)} \frac{A_{1,1}^2}{2} \sin 2k_y y \tag{11}$$

which shows that the spatial harmonic in the direction of the gradient is generated through this nonlinearity. We call it a new mode

$$\left(\frac{\tilde{n}}{n_0}\right)_{2,0} = A_{2,0} \sin 2k_y y. \tag{12}$$

One can write the continuity Eq. (6) in a form

$$\frac{\partial \tilde{n}}{\partial t} = \gamma\, \tilde{n} - \frac{c\underline{E}_\perp \times \hat{z}}{B_0} \cdot \nabla_\perp \tilde{n} \tag{13}$$

where we have split $n = n_0 + \tilde{n}$, $n_0(y)$ is the ambient plasma density, $\tilde{n}$ is the perturbation, and γ is the net growth (or damping) of the modes. It can be readily verified that a perturbation of the form (9) grows with the growth rate given by expression (8) while the mode (12) simply damps out due to diffusion. Following the procedure outlined in *Rognlien and Weinstock* [1974] and *Chaturvedi and Ossakow* [1977; 1979], we substitute for the density perturbation, in Eq. (13),

$$\frac{\tilde{n}}{n_o} = A_{1,1} \sin (k_x x + k_z z - \omega t) \cos k_y y + A_{2,0} \sin 2k_y y \tag{14}$$

The associated potential perturbation then can be computed from the linearized Eq. (7) and substituted in Eq. (13).

One then obtains two coupled equations for the mode amplitudes,

$$\frac{\partial A_{1,1}}{\partial t} = \gamma_{1,1} A_{1,1} - \alpha\, A_{1,1} A_{2,0} \tag{15}$$

and

$$\frac{\partial A_{2,0}}{\partial t} = -\gamma_{2,0}\, A_{2,0} + \frac{\alpha}{2} A_{1,1}^2 \tag{16}$$

where

$$\alpha = \frac{\dfrac{k_x k_y}{k_\perp^2}\,\dfrac{k_z v_d}{\Omega_e}}{\left(\dfrac{k_z^2}{k_\perp^2} + \dfrac{\nu_{ei}\nu_{in}}{\Omega_e \Omega_i}\right)} \nu_{ei}$$

is the coupling coefficient. One notes that in Eqs. (15) and (16), the nonlinearity causes a nonlinear damping to be introduced for the linearly growing mode $A_{1,1}$ through its interaction with mode $A_{2,0}$; whereas the linearly damped mode $A_{2,0}$ gets driven on account of the nonlinear interaction of mode $A_{1,1}$ with itself. In an asymptotic steady state, one can put $\partial A_{1,1}/\partial t = \partial A_{2,0}/\partial t = 0$. Then (15) and (16) yield,

$$A_{2,0} = \frac{\gamma_{1,1}}{\alpha} \simeq \frac{1}{k_y L} \tag{17}$$

and

$$A_{1,1} = \left\{ \frac{2\gamma_{2,0}}{\alpha} A_{2,0} \right\}^{1/2} \tag{18}$$

A rough estimate may be made of the saturated amplitudes. If $\gamma_{1,1} > \gamma_{2,0}$, then one notes that $A_{2,0} > A_{1,1}$. Then for $L \sim 50$ km and $\lambda_y \sim 10$ km, one obtains for $A_{2,0} \sim 3.2\%$. It may be mentioned here that in a recent in-situ rocket measurement of auroral F-region irregularities, a percentage fluctuation amplitude of this order was measured for these scale sizes (M.C. Kelley, private communication, 1979). It should also be noted that the linear theory of the current convective instability favors a wave vector in the east-west direction (see Eq. (8)), i.e., perpendicular to the northward density gradient. However, with $A_{2,0} > A_{1,1}$, the nonlinear saturated state of the instability would make the k vector appear to be in the y (northward) direction, i.e., the final spectrum would be dominated by the north-south modes. This is in agreement with the observations of *Rino et al.* [1978] who observed the sheet-like scintillation causing F region irregularities to be L shell aligned. Furthermore, the simulations of *Keskinen et al.* [1979] have shown that the collisional Rayleigh-Taylor instability in equatorial spread F saturates by the two dimensional mode coupling.

In summary, we have shown that a nonlinear two-dimensional mode coupling can stabilize the current convective instability in the auroral F-region ionosphere during diffuse auroral conditions. The saturated fluctuation amplitudes of these large scalesize structures (~10 km) could be on the order of a few percent of the background values.

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